

# Pseudoscalar condensation induced by chiral anomaly and vorticity for massive fermions

Ren-hong Fang,<sup>1</sup> Jin-yi Pang,<sup>2</sup> Qun Wang,<sup>1</sup> and Xin-nian Wang<sup>3,4</sup>

<sup>1</sup>*Interdisciplinary Center for Theoretical Study and Department of Modern Physics,  
University of Science and Technology of China, Hefei, Anhui 230026, China*

<sup>2</sup>*Helmholtz-Institut fuer Strahlen- und Kernphysik (Theorie) and Bethe  
Center for Theoretical Physics, Universitaet Bonn, D-53115 Bonn, Germany*

<sup>3</sup>*Key Laboratory of Quark and Lepton Physics (MOE) and Institute of Particle Physics,  
Central China Normal University, Wuhan, 430079, China*

<sup>4</sup>*Nuclear Science Division, MS 70R0319, Lawrence Berkeley National Laboratory, Berkeley, California 94720*

We derive the pseudoscalar condensate induced by anomaly and vorticity from the Wigner function for massive fermions in homogeneous electromagnetic fields. It has an anomaly term and a force-vorticity coupling term. As a mass effect, the pseudoscalar condensate is linearly proportional to the fermion mass in small mass expansion. By a generalization to two-flavor and three-flavor cases, the neutral pion and eta meson condensates are calculated from the Wigner function and have anomaly parts as well as force-vorticity parts, in which the anomaly part of the neutral pion condensate is consistent to the previous result. We also discuss about possible observables of the condensates in heavy ion collisions such as collective flows of neutral pions and eta mesons which may be influenced by the electromagnetic field and vorticity profiles.

## I. INTRODUCTION

The chiral or axial anomaly is the anomalous nonconservation of a chiral or axial current of fermions arising from the quantum tunneling effect between vacuums of gauge fields with different winding numbers. Chiral anomaly is also called Adler-Bell-Jackiw (ABJ) anomaly after the names of three founders [1, 2]. In quantum electrodynamics the anomalous nonconservation of the chiral or axial current can be written as

$$\partial_\mu j_5^\mu = -2mP - \frac{Q^2}{8\pi^2} F_{\mu\nu} \tilde{F}^{\mu\nu}, \quad (1)$$

where  $m$  and  $Q$  are the fermion mass and charge respectively,  $F_{\mu\nu}$  is the strength tensor of electromagnetic field with  $\tilde{F}^{\rho\lambda} = \frac{1}{2}\epsilon^{\rho\lambda\mu\nu}F_{\mu\nu}$  being its dual, chiral or axial current is defined by  $j_5^\mu = \bar{\psi}\gamma^\mu\gamma_5\psi$ , the pseudoscalar is defined by  $P = -i\bar{\psi}\gamma_5\psi$ , where  $\psi$  and  $\bar{\psi}$  are fermionic fields,  $\gamma^\mu$  ( $\mu = 0, 1, 2, 3$ ) are Dirac matrices and  $\gamma^5 = i\gamma^0\gamma^1\gamma^2\gamma^3$  is the chiral matrix. The most successful test of chiral anomaly is in the decay of a neutral pion into two photons, which had been a puzzle for some time in 1960s whose solution led to the discovery of the ABJ anomaly. For neutral pions, one can define the chiral current as  $j_{5,\pi}^\mu = \bar{\psi}\gamma^\mu\gamma_5(\sigma_3/2)\psi$  and pseudoscalar as  $P_\pi = -i\bar{\psi}\gamma_5(\sigma_3/2)\psi$ , where  $\psi = (u, d)^T$  and  $\bar{\psi} = (\bar{u}, \bar{d})$  are quark fields of two flavors and  $\sigma_3 = \text{diag}(1, -1)$  is the third Pauli matrix. The anomaly equation (1) now becomes

$$\partial_\mu j_{5,\pi}^\mu = f_\pi m_\pi^2 \phi_\pi - \frac{Q_e^2}{16\pi^2} F_{\mu\nu} \tilde{F}^{\mu\nu}, \quad (2)$$

where  $Q_e$  is the absolute value of the electron's charge. Here we have used the PCAC (partially conserved axial current) hypothesis [3] to relate the pseudoscalar  $P_\pi$  to the neutral pion field  $\phi_\pi$ ,  $2m_q P_\pi = -f_\pi m_\pi^2 \phi_\pi$ , where  $f_\pi$  is the pion decay constant and  $m_q$  and  $m_\pi$  are the quark and pion mass respectively. In the chiral limit with zero quark mass we have  $m_\pi = 0$  indicating pions as Goldstone bosons.

The chiral magnetic effect (CME) is an effect closely related to the chiral anomaly [4–7]. It is about the generation of an electric current along the magnetic field resulting from an imbalance of the population of chiral fermions. Another accompanying effect is the vortical effect in which an electric current is induced by the vorticity in a system of charged particles [8–10]. For chiral fermions it is called the chiral vortical effect (CVE) [11–13]. It has been demonstrated that the electric current from CME and CVE must coexist in order to guarantee the second law of thermodynamics in a chiral fluid [11, 14]. The CVE can be regarded as a quantum effect in hydrodynamics related to chiral anomaly.

The CME, CVE and other related effects such as chiral magnetic wave [15, 16] have been extensively studied in the quark-gluon plasma produced in high-energy heavy-ion collisions in which very strong magnetic fields [5, 17–25] and huge global angular momenta [26–29] are produced in non-central collisions. The charge separation effect observed in STAR [30, 31] and ALICE [32] experiments are consistent to the CME prediction. But there were debates that the charge separation might come from cluster particle correlations together with a reasonable range of cluster anisotropy

in nonperipheral collisions [33], so it is not conclusive that the charge separation effect would be the evidence of the CME. The charge asymmetry dependence of pion elliptic flow was observed in heavy-ion collisions by STAR and is considered as the possible consequences of the chiral magnetic wave [34]. The CME has recently been confirmed to exist in materials such as Dirac and Weyl semi-metals [35–37]. Recently STAR collaboration has measured nonvanishing hyperon polarization in the beam energy scan program [38]. This is a piece of evidence for local polarization effect from vorticity in collisions at lower energy and was first predicted in Ref. [39].

Quantum kinetic theory in terms of Wigner function [40–43] is a useful tool to study the CME, CVE and other related effects [13, 44–46]. The axial vector component of the Wigner function for massless fermions can be generalized to massive fermions and gives their phase-space density of the spin vector. The spin vector arises from nonzero fermion mass [47]. Therefore one can calculate the polarization of massive fermions from the axial vector component [46]. The polarization density is found to be proportional to the local vorticity  $\omega$  as well as the magnetic field. The polarization per particle for fermions is always smaller than that for anti-fermions as the result of more Pauli blocking effect for fermions than antifermions. This is consistent to the STAR’s preliminary result on the  $\Lambda$  polarization [38].

In this paper we give another important feature of massive fermions from the axial vector component of the Wigner function, namely, the thermal average of pseudoscalar quantity  $P$  in Eq. (1). In low energy QCD, it is proportional to the pion field from PCAC hypothesis, so we call it the pseudoscalar condensate. We will show that such a pseudoscalar condensate is induced by anomaly and force-vorticity coupling and depends on the fermion mass, fermion chemical potential and temperature. It is a mass effect in a plasma of fermions: for massless fermions the pseudoscalar is vanishing. We note that there were many studies on pseudoscalar condensates such as pion and eta meson condensates in nuclear and quark matter in a variety of hot and dense environments [48–52] which are different from the condensates on which we focus in this paper.

The paper is organized as follows. In Section II, we summarize the properties of the Wigner function for massless fermions in electromagnetic fields. In Section III we present the equation for pseudoscalar and axial vector component of the Wigner function for massive fermions, the Wigner function counterpart of Eq. (1). In Section IV, we analyze the axial vector component at leading order from which the polarization vector can be obtained. We also derive the polarization vector of a fermion in the lab frame with its 3-momentum  $\mathbf{p}$  in the fluid cell’s comoving frame and the polarization vector  $\mathbf{n}$  in the particle’s rest frame. This is useful to connect the experimental observable to the theoretical prediction. In Section V we derive nonconservation of chiral current with anomaly and fermion mass by taking space-time divergence of the chiral current derived from the axial vector component. We also calculate condensates of neutral pions and eta mesons. Then we derive the pseudoscalar condensate induced by anomaly and vorticity for massive fermions. The summary is made in the last section.

We adopt the same sign conventions for fermion charge  $Q$  as in Refs. [13, 42, 44, 45], and the same sign convention for the axial vector component  $\mathcal{A}^\mu \sim \langle \bar{\psi} \gamma^\mu \gamma^5 \psi \rangle$  as in Refs. [13, 44, 45] but different sign convention from Refs. [42].

## II. WIGNER FUNCTION FOR MASSLESS FERMIONS IN ELECTROMAGNETIC FIELDS

The gauge invariant Wigner function is the quantum mechanical analogue of a classical phase-space distribution. In a background electromagnetic field, the Wigner function  $W_{\alpha\beta}(x, p)$  is defined by

$$W_{\alpha\beta}(x, p) = \int \frac{d^4 y}{(2\pi)^4} e^{-ip \cdot y} \left\langle \bar{\psi}_\beta(x + \frac{1}{2}y) PU(G, x + \frac{1}{2}y, x - \frac{1}{2}y) \psi_\alpha(x - \frac{1}{2}y) \right\rangle, \quad (3)$$

where  $\psi_\alpha$  and  $\bar{\psi}_\beta$  are fermionic quantum fields with Dirac indices  $\alpha$  and  $\beta$ ,  $\langle \hat{O} \rangle$  denotes the grand canonical ensemble average of normal ordered operator,  $x = (x_0, \mathbf{x})$  and  $p = (p_0, \mathbf{p})$  are time-space and energy-momentum 4-vectors respectively, and the gauge link  $PU(G, x_1, x_2)$  is to ensure the gauge invariance of the Wigner function where  $G^\mu$  is the gauge potential of the classical electromagnetic field and  $P$  denotes the path-ordered product. The  $4 \times 4$  matrix  $W_{\alpha\beta}(x, p)$  can be decomposed by 16 independent generators of Clifford algebra, namely,  $1, \gamma^5, \gamma^\mu, \gamma^5 \gamma^\mu, \sigma^{\mu\nu} \equiv \frac{i}{2}[\gamma^\mu, \gamma^\nu]$ , into the scalar, pseudoscalar, vector, axial vector and tensor components, respectively. From the Dirac equation for the fermionic field, one can derive the equation for  $W_{\alpha\beta}(x, p)$  which leads to a set of coupled equations for all components. Finding a general solution to the Wigner function is very difficult. However it is much simplified for massless fermions for which the set of equations for the vector and axial vector components are decoupled from the rest components. Assuming that the electromagnetic field is homogeneous and weak and is in the same order as the space-time derivative, one can solve the vector and axial vector components perturbatively. To the linear order in the field strength and vorticity, one can obtain the vector and axial vector components [13], which give the charge current  $j^\mu$  and chiral (axial) charge current  $j_5^\mu$  by integration over four-momenta:  $j^\mu = nu^\mu + \xi\omega^\mu + \xi_B B^\mu$  and  $j_5^\mu = nu^\mu + \xi_5\omega^\mu + \xi_{5B} B^\mu$ . Here  $\omega^\mu = \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}u_\nu\partial_\rho u_\sigma$  is the vorticity vector,  $B^\mu = \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}u_\nu F_{\rho\sigma}$  is the magnetic field four-vector with the fluid four-velocity  $u_\nu$ ,  $n$  and  $n_5$  are charge and chiral (axial) charge density respectively.

In the charge current  $j^\mu$  one obtains the CME and CVE coefficients  $\xi = \mu\mu_5/\pi^2$  and  $\xi_B = Q\mu_5/(2\pi^2)$  respectively, where  $\mu$  and  $\mu_5$  are chemical potentials for the charge and chiral (axial) charge respectively. One can also derive the coefficients of vorticity and magnetic field in  $j_5^\mu$ :  $\xi_5 = T^2/6 + (\mu^2 + \mu_5^2)/(2\pi^2)$  and  $\xi_{5B} = Q\mu/(2\pi^2)$ . The conservation and anomalous nonconservation laws for the charge and chiral (axial) charge current respectively can be verified,  $\partial_\mu j^\mu = 0$  and  $\partial_\mu j_5^\mu = -[Q^2/(8\pi^2)]F_{\mu\nu}\tilde{F}^{\mu\nu}$ . The covariant chiral kinetic equation which is related to the Berry phase in 4-dimension can also be derived from the first order solution to the vector and axial vector components of the Wigner function [44].

### III. EQUATION FOR PSEUDOSCALAR AND AXIAL VECTOR COMPONENT OF WIGNER FUNCTION

In this section we look at the equation for pseudoscalar and axial vector component of the Wigner function. From Dirac equation for the fermionic field, one can derive the equation for the Wigner function in (3) in a constant electromagnetic field,

$$\left[ \gamma_\mu \left( p^\mu + i\hbar \frac{1}{2} \nabla^\mu \right) - m \right] W(x, p) = 0, \quad (4)$$

where the phase-space derivative is defined by  $\nabla^\mu = \partial_x^\mu - QF^{\mu\nu}\partial_{p,\nu}$  and we suppressed Dirac indices of the Wigner function. The Wigner function as a  $4 \times 4$  matrix in Dirac space can be decomposed into the scalar, pseudoscalar, vector, axial vector and tensor components as

$$W = \frac{1}{4} \left[ \mathcal{F} + i\gamma^5 \mathcal{P} + \gamma^\mu \mathcal{V}_\mu + \gamma^5 \gamma^\mu \mathcal{A}_\mu + \frac{1}{2} \sigma^{\mu\nu} \mathcal{S}_{\mu\nu} \right]. \quad (5)$$

The components in the decomposition (5) can be obtained by projection of corresponding Dirac matrices on the Wigner function and taking traces. Eq. (4) for the Wigner function can be converted to a set of coupled equations for all components. There is an equation that relates the pseudoscalar to the axial vector component which is of special interest,

$$\hbar \nabla^\mu \mathcal{A}_\mu = -2m \mathcal{P}. \quad (6)$$

An interesting observation of the above equation is that the pseudoscalar component  $\mathcal{P}$  is of quantum origin since it is proportional to the Planck constant  $\hbar$ . We note that Eq. (6) is nothing but the Wigner function counterpart of Eq. (1). We will show that the integration of Eq. (6) over 4-momentum gives Eq. (1).

### IV. AXIAL VECTOR COMPONENT AT LEADING ORDER AND POLARIZATION VECTOR

At leading (zero-th) order of electromagnetic interaction, the gauge link in the Wigner function in Eq. (3) can be set to 1, we denote the Wigner function at this order as  $W_{(0)}$ . We can expand fermionic fields in momentum space with creation and destruction operators, which we insert into Eq. (3). After taking ensemble average of normal ordered operators, we obtain

$$\begin{aligned} W_{(0)}(x, p) = & \frac{1}{(2\pi)^3} \delta(p^2 - m^2) \left\{ \theta(p^0) \sum_s f_{\text{FD}}(E_p - \mu_s) u(\mathbf{p}, s) \bar{u}(\mathbf{p}, s) \right. \\ & \left. - \theta(-p^0) \sum_s f_{\text{FD}}(E_p + \mu_s) v(-\mathbf{p}, s) \bar{v}(-\mathbf{p}, s) \right\}, \end{aligned} \quad (7)$$

where  $u(\mathbf{p}, s)$  and  $v(-\mathbf{p}, s)$  are Dirac spinors of positive and negative energy respectively,  $s = \pm$  denote the spin state parallel or anti-parallel to the spin quantization direction  $\mathbf{n}$  in the rest frame of the particle. We have also used  $\langle a^\dagger(\mathbf{p}, s) a(\mathbf{p}, s) \rangle = f_{\text{FD}}(E_p - \mu_s)$  and  $\langle b^\dagger(-\mathbf{p}, s) b(-\mathbf{p}, s) \rangle = f_{\text{FD}}(E_p + \mu_s)$  with the Fermi-Dirac distribution defined by  $f_{\text{FD}}(x) = 1/(e^{\beta x} + 1)$  ( $\beta \equiv 1/T$ ,  $T$  is temperature) and  $\mu_s$  is the chemical potential for fermions in the spin state  $s$ .

The axial vector component at the leading order is given by

$$\begin{aligned} \mathcal{A}_{(0)}^\mu &= \text{Tr}[\gamma^\mu \gamma^5 W_{(0)}] \\ &= m [\theta(p_0) n^\mu(\mathbf{p}, \mathbf{n}) - \theta(-p_0) n^\mu(-\mathbf{p}, -\mathbf{n})] \delta(p^2 - m^2) A, \end{aligned} \quad (8)$$

where  $A$  is defined by

$$A \equiv \frac{2}{(2\pi)^3} \sum_s s [\theta(p^0) f_{\text{FD}}(p_0 - \mu_s) + \theta(-p^0) f_{\text{FD}}(-p_0 + \mu_s)], \quad (9)$$

and we have used  $\bar{u}(\mathbf{p}, s) \gamma^\mu \gamma^5 u(\mathbf{p}, s) = 2ms n^\mu(\mathbf{p}, \mathbf{n})$  and  $\bar{v}(-\mathbf{p}, s) \gamma^\mu \gamma^5 v(-\mathbf{p}, s) = 2ms n^\mu(-\mathbf{p}, -\mathbf{n})$  with  $n^\mu(\mathbf{p}, \mathbf{n})$  given by

$$n^\mu(\mathbf{p}, \mathbf{n}) = \Lambda^\mu_\nu(-\mathbf{v}_p) n^\nu(\mathbf{0}, \mathbf{n}) = \left( \frac{\mathbf{n} \cdot \mathbf{p}}{m}, \mathbf{n} + \frac{(\mathbf{n} \cdot \mathbf{p}) \mathbf{p}}{m(m + E_p)} \right). \quad (10)$$

Here  $\Lambda^\mu_\nu(-\mathbf{v}_p)$  is the Lorentz transformation for  $\mathbf{v}_p = \mathbf{p}/E_p$  and  $n^\nu(\mathbf{0}, \mathbf{n}) = (0, \mathbf{n})$  is the 4-vector of the spin quantization direction in the rest frame of the fermion. One can check that  $n^\mu(\mathbf{p}, \mathbf{n})$  satisfies  $n^2 = -1$  and  $n \cdot p = 0$ , so it behaves like a spin 4-vector up to a factor of  $1/2$ . For Pauli spinors  $\chi_s$  and  $\chi_{s'}$  in  $u(\mathbf{p}, s)$  and  $v(-\mathbf{p}, s')$  respectively, we have  $\chi_s^\dagger \boldsymbol{\sigma} \chi_s = s\mathbf{n}$  and  $\chi_{s'}^\dagger \boldsymbol{\sigma} \chi_{s'} = -s'\mathbf{n}$ . We can take the massless limit by setting  $\mathbf{n} = \hat{\mathbf{p}}$ , then we have  $mn^\mu(\mathbf{p}, \mathbf{n}) \rightarrow (|\mathbf{p}|, \mathbf{p})$  and  $mn^\mu(-\mathbf{p}, -\mathbf{n}) \rightarrow (|\mathbf{p}|, -\mathbf{p})$ . This way we can recover the previous result of the axial vector component for massless fermions [13, 44] where  $s = \pm$  denote the right-handed and left-handed fermions.

We implied that  $n^\mu(\mathbf{p}, \mathbf{n})$  in (10) is the form in the comoving or local rest frame of a fluid cell. We can boost all quantities in Eq. (8) to the lab frame in which the fluid cell is moving with a 4-velocity  $u^\alpha = \gamma(1, \mathbf{v})$  where  $\gamma = 1/\sqrt{1 - v^2}$  is the Lorentz factor. We boost  $n^\mu(\mathbf{p}, \mathbf{n})$  and  $n^\mu(-\mathbf{p}, -\mathbf{n})$  to the lab frame as

$$\begin{aligned} n_{\text{lab}}^{\prime\alpha}(\mathbf{p}, \mathbf{n}) &= \Lambda^\alpha_\beta(-\mathbf{v}) n^\beta(\mathbf{p}, \mathbf{n}), \\ n_{\text{lab}}^{\prime 0}(\mathbf{p}, \mathbf{n}) &= \gamma \left[ \frac{\mathbf{n} \cdot \mathbf{p}}{m} + \mathbf{v} \cdot \mathbf{n} + \frac{(\mathbf{n} \cdot \mathbf{p})(\mathbf{v} \cdot \mathbf{p})}{m(m + E_p)} \right], \\ \mathbf{n}'_{\text{lab}}(\mathbf{p}, \mathbf{n}) &= \mathbf{n} + \frac{(\mathbf{n} \cdot \mathbf{p}) \mathbf{p}}{m(m + E_p)} + \frac{\gamma - 1}{v^2} \mathbf{v} \left[ \mathbf{n} \cdot \mathbf{v} + \frac{(\mathbf{n} \cdot \mathbf{p})(\mathbf{p} \cdot \mathbf{v})}{m(m + E_p)} \right] + \frac{\gamma}{m} \mathbf{v}(\mathbf{n} \cdot \mathbf{p}). \end{aligned} \quad (11)$$

Note that  $\mathbf{n}$  is the spin quantization direction in the fermion's rest frame and  $\mathbf{p}$  is the fermion's momentum in the local rest frame of the fluid cell. The fermion momentum in the lab frame is  $p'^\alpha = \Lambda^\alpha_\beta(-\mathbf{v}) p^\beta$  with  $p^\beta = (E_p, \mathbf{p})$  or explicitly

$$\begin{aligned} p'^0 &= \gamma[E_p + (\mathbf{p} \cdot \mathbf{v})], \\ \mathbf{p}' &= \mathbf{p} + \frac{\gamma - 1}{v^2} \mathbf{v}(\mathbf{p} \cdot \mathbf{v}) + \gamma \mathbf{v} E_p. \end{aligned} \quad (12)$$

One can check  $p' \cdot u = p^0 = E_p$  and  $n' \cdot p' = 0$  in the lab frame. So Eq. (8) and (9) can be written in the lab frame by making the replacement  $n^\alpha(\mathbf{p}, \mathbf{n}) \rightarrow n_{\text{lab}}^{\prime\alpha}(\mathbf{p}, \mathbf{n})$ , also we have  $p_0 = p' \cdot u$  and  $p'^2 = p^2$  in the two formula.

We note that Eq. (11) is the polarization vector in the lab frame (the fluid cell moves with a velocity  $\mathbf{v}$ ) of a fermion with the 3-momentum  $\mathbf{p}$  in the fluid cell's comoving frame and the polarization vector  $\mathbf{n}$  in the particle's rest frame. After taking integration of  $n_{\text{lab}}^{\prime\alpha}(\mathbf{p}, \mathbf{n})$  over  $\mathbf{p}$  which follows the Fermi-Dirac distribution, one can obtain the thermal average of the polarization vector in the lab frame. The  $\Lambda$  polarization can be measured in experiments by its decay to proton and pion in its rest frame. It is observed that  $\Lambda$  is polarized along the global angular momentum in the beam energy scan program at RHIC [38]. From Eq. (11) one can calculate the  $\Lambda$  polarization along a fixed direction in the lab frame, e.g. the direction of the global angular momentum, and compare with data.

In this paper, we assume that the chemical potential does not depend on the spin state  $\mu_s = \mu$  for  $s = \pm$ , so  $A = 0$  and then  $\mathcal{A}_{(0)}^\mu$  is vanishing. The nonvanishing contribution comes from the axial vector component at the next-to-leading or first order even with spin-independent chemical potentials.

## V. CHIRAL CURRENT NONCONSERVATION LAW AND PSEUDOSCALAR CONDENSATION

In this section we will derive nonconservation law of chiral current of massive fermions by taking space-time divergence of the chiral current derived from the axial vector component. Then we can derive the pseudoscalar condensate induced by anomaly and vorticity.

As we have discussed in the last section that the contribution should come from the axial vector component at the next-to-leading or first order. We use the following form for the axial vector component for massive fermions by generalizing the solution for massless fermions [13, 44–46],

$$\mathcal{A}_{(1)}^\alpha(x, p) = -\frac{1}{2} \hbar \beta \tilde{\Omega}^{\alpha\sigma} p_\sigma \frac{dV}{d(\beta p_0)} \delta(p^2 - m^2) - Q \hbar \tilde{F}^{\alpha\lambda} p_\lambda V \frac{\delta(p^2 - m^2)}{p^2 - m^2}, \quad (13)$$

where  $p_0 \equiv u \cdot p$ ,  $\tilde{\Omega}^{\rho\sigma} = \frac{1}{2}\epsilon^{\rho\sigma\mu\nu}\partial_\mu u_\nu$ , and  $V$  is associated with the vector component and given by

$$V \equiv \frac{4}{(2\pi)^3} [\theta(p^0)f_{\text{FD}}(p_0 - \mu) + \theta(-p^0)f_{\text{FD}}(-p_0 + \mu)]. \quad (14)$$

where we have taken  $\mu_s = \mu$  for  $s = \pm$ . Also we have assumed in Eq. (13) that  $\beta = 1/T$  is a constant. The chiral current can be obtained by integrating over 4-momentum of the axial vector component,  $j_5^\mu = \int d^4p \mathcal{A}_{(1)}^\mu(x, p)$ , whose space-time divergence is given by

$$\begin{aligned} \partial_\mu j_5^\mu &= \int d^4p \partial_\mu \mathcal{A}_{(1)}^\mu \\ &= -\frac{1}{2}\hbar\beta\tilde{\Omega}^{\mu\sigma} \int d^4p p_\sigma \partial_\mu \left[ \frac{dV}{d(\beta p_0)} \right] \delta(p^2 - m^2) \\ &\quad - \hbar Q \tilde{F}^{\mu\lambda} \int d^4p p_\lambda (\partial_\mu V) \frac{\delta(p^2 - m^2)}{p^2 - m^2}. \end{aligned} \quad (15)$$

where we have assumed that  $\tilde{F}^{\mu\lambda}$  does not depend on space-time, and used  $u_\sigma \partial_\mu \tilde{\Omega}^{\mu\sigma} = 0$  due to static equilibrium conditions [13]. When taking space-time divergence, we neglect the space-time derivative of  $\theta(p_0)$  and  $\theta(-p_0)$  in  $V$  which would be vanishing when contacting the mass-shell condition. The two terms in the last equality of Eq. (15) are evaluated in Appendix A and can be grouped into a  $E \cdot B$  and  $E \cdot \omega$  term,

$$\begin{aligned} \partial_\mu j_5^\mu &= \frac{1}{2}\hbar Q^2 \beta^2 (E \cdot B) \int d^4p V''_{\beta\mu, \beta p_0} \delta(p^2 - m^2) \\ &\quad + \frac{1}{2}\hbar\beta Q (E \cdot \omega) \int d^4p \left[ \beta p_0 V''_{\beta p_0, \beta\mu} - \beta \frac{2\bar{p}^2}{3p_0} V''_{\beta p_0, \beta p_0} + \frac{2\bar{p}^2}{3p_0^2} V'_{\beta p_0} \right] \delta(p^2 - m^2), \end{aligned} \quad (16)$$

where  $\bar{p}^\alpha \equiv p^\alpha - (p \cdot u)u^\alpha$ . The integrals in Eq. (16) can be finally simplified into following forms

$$\partial_\mu j_5^\mu = -\frac{1}{2\pi^2}\hbar Q^2 (E \cdot B) C_1(\beta m, \beta\mu) - \frac{m^2}{2\pi^2}\hbar Q \beta (E \cdot \omega) C_2(\beta m, \beta\mu), \quad (17)$$

where the dimensionless functions  $C_1(\beta m, \beta\mu)$  and  $C_2(\beta m, \beta\mu)$  are defined by

$$\begin{aligned} C_1(\beta m, \beta\mu) &= \int_0^\infty dx \left[ \frac{e^{\sqrt{x^2 + (\beta m)^2} - \beta\mu}}{(e^{\sqrt{x^2 + (\beta m)^2} - \beta\mu} + 1)^2} + \frac{e^{\sqrt{x^2 + (\beta m)^2} + \beta\mu}}{(e^{\sqrt{x^2 + (\beta m)^2} + \beta\mu} + 1)^2} \right], \\ C_2(\beta m, \beta\mu) &= \int_0^\infty dx \frac{1}{\sqrt{x^2 + (\beta m)^2}} \left[ \frac{e^{\sqrt{x^2 + (\beta m)^2} - \beta\mu}}{(e^{\sqrt{x^2 + (\beta m)^2} - \beta\mu} + 1)^2} - \frac{e^{\sqrt{x^2 + (\beta m)^2} + \beta\mu}}{(e^{\sqrt{x^2 + (\beta m)^2} + \beta\mu} + 1)^2} \right]. \end{aligned} \quad (18)$$

We plot the functions  $C_1(\beta m, \beta\mu) - 1$  and  $C_2(\beta m, \beta\mu)$  in Fig. 1. One can verify asymptotic values  $C_1(\beta m, \beta\mu) - 1 \sim O[(\beta m)^2]$  and  $C_2(\beta m, \beta\mu) \sim O(1)$  for  $\beta m \rightarrow 0$ , see Eq. (B19) in Appendix B.

On the other hand, taking an integration over 4-momentum of Eq. (6) leads to

$$-2mP = \hbar \partial_\mu j_5^\mu + \frac{1}{2\pi^2}\hbar^2 Q^2 (E \cdot B), \quad (19)$$

where we have used  $P(x) = \int d^4p \mathcal{P}(x, p)$  and the inetgral

$$Q F_{\mu\nu} \int d^4p \partial_p^\nu \mathcal{A}^\mu(x, p) = -\frac{1}{2\pi^2}\hbar Q^2 (E \cdot B), \quad (20)$$

whose derivation is given in Appendix C. Inserting Eq. (17) into Eq. (19) we obtain the pseudoscalar condensate

$$P = \frac{1}{4\pi^2}\hbar^2 Q^2 (E \cdot B) \frac{1}{m} [C_1(\beta m, \beta\mu) - 1] + \frac{1}{4\pi^2}\hbar^2 Q (E \cdot \omega) \beta m C_2(\beta m, \beta\mu). \quad (21)$$

From the small mass behavior  $C_1(\beta m, \beta\mu) - 1 \sim (\beta m)^2$  and  $C_2(\beta m, \beta\mu) \sim O(1)$ , the pseudoscalar is proportional to the fermion mass. This feature is similar to the PCAC hypothesis where the pseudoscalar is proportional to  $m_\pi^2/m_q$ . We see in  $C_2(\beta m, \beta\mu)$  that there is a sign difference between fermion and antifermion terms, so the  $QE \cdot \omega$  term can also be regarded as the *force-vorticity* coupling term. Furthermore the  $E \cdot \omega$  term is vanishing for  $\beta\mu = 0$ .

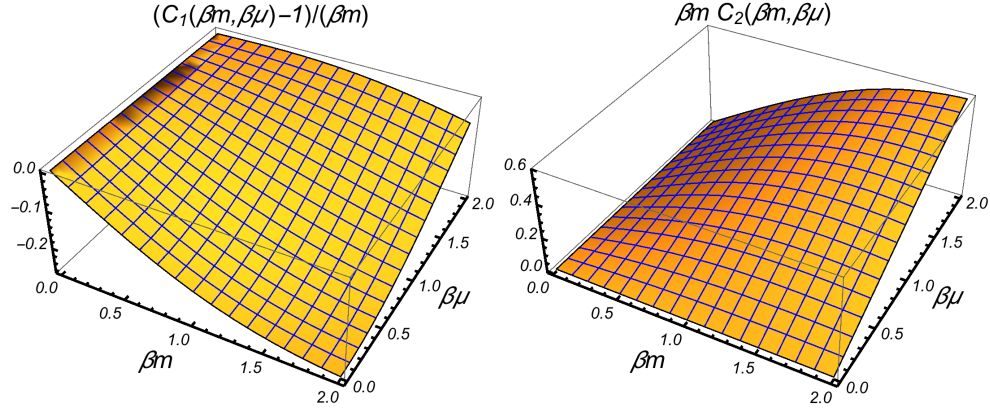


Figure 1:  $[C_1(\beta m, \beta \mu) - 1]/(\beta m)$  and  $(\beta m)C_2(\beta m, \beta \mu)$  as functions of  $\beta m$  and  $\beta \mu$ . They are all proportional to  $m$  in the chiral limit.

The above discussions are about single fermion species. Let us consider a quark plasma with two quark flavors  $u$  and  $d$ . For  $u$  or  $d$  quarks, Eqs. (17,19) become

$$\begin{aligned} \partial_\mu j_{5,a}^\mu &= -\frac{N_C}{2\pi^2} \hbar Q_a^2 (E \cdot B) C_1(\beta m_a, \beta \mu_a) - \frac{N_C m_a^2}{2\pi^2} \hbar Q_a \beta (E \cdot \omega) C_2(\beta m_a, \beta \mu_a), \\ -2m_a P_a &= \hbar \partial_\mu j_{5,a}^\mu + \frac{N_C}{2\pi^2} \hbar^2 Q_a^2 (E \cdot B), \end{aligned} \quad (22)$$

where  $N_C = 3$  is the number of color,  $a = u, d$ , and  $P_a = -i\bar{a}\gamma_5 a$ ,  $j_{5,a}^\mu = \bar{a}\gamma^\mu\gamma_5 a$ . We know from the introduction that the pseudoscalar and chiral current for neutral pions are given by

$$\begin{aligned} P_\pi &= -i\bar{\psi}\gamma_5(\sigma_3/2)\psi = \frac{1}{2}(P_u - P_d), \\ j_{5,\pi}^\mu &= \bar{\psi}\gamma^\mu\gamma_5(\sigma_3/2)\psi = \frac{1}{2}(j_{5,u}^\mu - j_{5,d}^\mu). \end{aligned} \quad (23)$$

In the first case we assume that the masses and chemical potentials of  $u$  and  $d$  quarks are the same,  $m_u = m_d = m_q$  and  $\mu_u = \mu_d = \mu_q$ . Then the pion condensate is

$$\begin{aligned} P_\pi &= \frac{N_C}{8\pi^2} \hbar^2 (Q_u^2 - Q_d^2) (E \cdot B) \frac{1}{m_q} [C_1(\beta m_q, \beta \mu_q) - 1] \\ &\quad + \frac{N_C}{8\pi^2} \hbar^2 (Q_u - Q_d) (E \cdot \omega) \beta m_q C_2(\beta m_q, \beta \mu_q) \\ &= \frac{N_C}{24\pi^2} \hbar^2 Q_e^2 (E \cdot B) \frac{1}{m_q} [C_1(\beta m_q, \beta \mu_q) - 1] + \frac{N_C}{8\pi^2} \hbar^2 Q_e (E \cdot \omega) \beta m_q C_2(\beta m_q, \beta \mu_q) \end{aligned} \quad (24)$$

In the second case we assume that the masses of  $u$  and  $d$  quarks are the same,  $m_u = m_d = m_q$ , but the chemical potentials are different, the pion condensate is

$$\begin{aligned} P_\pi &= \frac{N_C}{8\pi^2} \hbar^2 (E \cdot B) \frac{1}{m_q} [Q_u^2 C_1(\beta m_q, \beta \mu_u) - Q_d^2 C_1(\beta m_q, \beta \mu_d) - (Q_u^2 - Q_d^2)] \\ &\quad + \frac{N_C}{8\pi^2} \hbar^2 (E \cdot \omega) \beta m_q [Q_u C_2(\beta m_q, \beta \mu_u) - Q_d C_2(\beta m_q, \beta \mu_d)] \end{aligned} \quad (25)$$

In vacuum at zero temperature and chemical potentials, both functions  $C_1$  and  $C_2$  are vanishing and we obtain from Eqs. (24,25)

$$P_\pi^{\text{vac}} = -\frac{\hbar^2}{8\pi^2 m_q} Q_e^2 (E \cdot B), \quad (26)$$

which is consistent with the result derived in the NJL model and chiral perturbation theory [53]. We see that the pion condensate in vacuum comes out quite natural as a mass effect of charged fermions and is not subject to any

additional constraint such as a critical value for  $|\mathbf{E} \cdot \mathbf{B}|$  in [53]. Recently the chiral current and pseudoscalar condensate for massive fermions have also been calculated in a holographic model in finite density and magnetic field [54].

We can also generalize to quark matter with three quark flavors  $u$ ,  $d$  and  $s$ . So Eq. (22) applies to  $a = u, d, s$ . Now we look at the pseudoscalar and chiral current for the flavor octet  $\eta_8$  and singlet  $\eta_1$  which can be defined by

$$\begin{aligned} P_{\eta_8} &= -i\bar{\psi}\gamma_5\frac{\lambda_8}{2}\psi = \frac{1}{2\sqrt{3}}(P_u + P_d - 2P_s), \\ P_{\eta_1} &= -i\bar{\psi}\gamma_5\frac{1}{\sqrt{3}}\mathbf{1}\psi = \frac{1}{\sqrt{3}}(P_u + P_d + P_s), \\ j_{5,\eta_8}^\mu &= \bar{\psi}\gamma^\mu\gamma_5\frac{\lambda_8}{2}\psi = \frac{1}{2\sqrt{3}}\left(j_{5,u}^\mu + j_{5,d}^\mu - 2j_{5,s}^\mu\right), \\ j_{5,\eta_1}^\mu &= \bar{\psi}\gamma^\mu\gamma_5\frac{1}{\sqrt{3}}\mathbf{1}\psi = \frac{1}{\sqrt{3}}\left(j_{5,u}^\mu + j_{5,d}^\mu + j_{5,s}^\mu\right). \end{aligned} \quad (27)$$

As an example for the  $\eta$  meson condensate, we consider the simplest case for quark masses and chemical potentials,  $m_u = m_d = m_s = m_q$  and  $\mu_u = \mu_d = \mu_s = \mu_q$ , which can be considered as the case of approaching the chiral limit. In this case, we have

$$\begin{aligned} P_{\eta_8} &= \frac{N_C}{8\sqrt{3}\pi^2}\hbar^2(Q_u^2 + Q_d^2 - 2Q_s^2)(E \cdot B)\frac{1}{m_q}[C_1(\beta m_q, \beta \mu_q) - 1] \\ &\quad + \frac{N_C}{8\sqrt{3}\pi^2}\hbar^2(Q_u + Q_d - 2Q_s)(E \cdot \omega)\beta m_q C_2(\beta m_q, \beta \mu_q) \\ &= \frac{N_C}{24\sqrt{3}\pi^2}\hbar^2 Q_e^2(E \cdot B)\frac{1}{m_q}[C_1(\beta m_q, \beta \mu_q) - 1] + \frac{N_C}{8\sqrt{3}\pi^2}\hbar^2 Q_e(E \cdot \omega)\beta m_q C_2(\beta m_q, \beta \mu_q) \\ P_{\eta_1} &= \frac{N_C}{4\sqrt{3}\pi^2}\hbar^2(Q_u^2 + Q_d^2 + Q_s^2)(E \cdot B)\frac{1}{m_q}[C_1(\beta m_q, \beta \mu_q) - 1] \\ &\quad + \frac{N_C}{4\sqrt{3}\pi^2}\hbar^2(Q_u + Q_d + Q_s)(E \cdot \omega)\beta m_q C_2(\beta m_q, \beta \mu_q) \\ &= \frac{N_C}{6\sqrt{3}\pi^2}\hbar^2 Q_e^2(E \cdot B)\frac{1}{m_q}[C_1(\beta m_q, \beta \mu_q) - 1]. \end{aligned} \quad (28)$$

The main difference between  $P_{\eta_8}$  and  $P_{\eta_1}$  in the above is that  $P_{\eta_1}$  does not have  $E \cdot \omega$  term due to the cancellation of electric charges of quarks with different flavors. Of course more realistic cases in heavy ion collisions should be like  $m_u = m_d = m_q \neq m_s$  and  $\mu_u = \mu_d = \mu_q \neq \mu_s$ , for which the calculation is straightforward. There are possible observables of  $\eta$  meson condensates in heavy ion collisions [50], but it is beyond the scope of the current paper and will be addressed in a future study.

The pseudoscalar condensate for charged fermions in Eq. (21) is our main result. It is a quite general formula. Such a pseudoscalar condensate is charge neutral and induced by anomaly and vorticity in thermal and dense environment, which is a natural consequence of nonconservation law of the chiral current in electromagnetic fields and is not subject to any additional constraints. For example, the neutral pion condensate is always in the forms of Eqs. (24,25) without further condition about the value of  $|\mathbf{E} \cdot \mathbf{B}|$  if anomaly and vorticity are there. It is also worth mentioning a new electric-field-vorticity coupling term in the condensate which has not been derived in literature to our knowledge. Such pion and eta meson condensates may have observables related to the electromagnetic field and vorticity in heavy ion collisions. For example they may have effects on the collective flows of neutral pions and eta mesons. In some sense they are similar to disoriented chiral condensates [55–57]. This is a topic that we will address in the future.

## VI. SUMMARY

We derive the pseudoscalar condensate induced by anomaly and vorticity from the Wigner function for massive fermions in homogeneous electromagnetic fields. The pseudoscalar component of the Wigner function is determined from the axial vector component by Eq. (6). Taking an integration over 4-momentum for Eq. (6) we obtain the anomalous nonconservation of the chiral current by the anomalous term and a product of mass and pseudoscalar. By directly calculating the space-time divergence of the chiral current, we can determine the pseudoscalar condensate which has an anomalous  $E \cdot B$  term and an  $E \cdot \omega$  term. The  $E \cdot \omega$  term can also be regarded as a force-vorticity coupling since there is a sign difference in its prefactor between the fermion and antifermion sector. The force-vorticity part

of the pseudoscalar condensate is the new term. As a mass effect, the pseudoscalar condensate is linearly proportional to the fermion mass when the mass is small. Such a pseudoscalar condensate is a general feature for a fluid of massive and charged fermions in a thermal and dense plasma with anomaly and vorticity. The neutral pion and eta meson condensates can also be derived from generalization of the single flavor to multi-flavor case, which depend on quark masses, quark chemical potentials and temperature. We can reproduce the previous result of the neutral pion condensate in vacuum induced by the anomaly, but our result also has a force-vorticity part which has not been derived in previous literature to our knowledge. There are possible observables of pseudoscalar condensates related to the electromagnetic field and vorticity in heavy ion collisions such as collective flows of neutral pions and eta mesons.

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### Appendix A: Derivation of Eq. (15)

Let us treat the vorticity term in Eq. (15) as,

$$\begin{aligned}
 I_\omega &= -\frac{1}{2}\hbar\beta\tilde{\Omega}^{\mu\sigma}\partial_\mu(\beta u_\rho)\int d^4p p_\sigma V''_{\beta p_0,\beta\mu}\delta(p^2-m^2) \\
 &\quad -\frac{1}{2}\hbar\beta\tilde{\Omega}^{\mu\sigma}\partial_\mu(\beta u_\rho)\int d^4p p_\sigma p^\rho V''_{\beta p_0,\beta p_0}\delta(p^2-m^2) \\
 &= \frac{1}{2}\hbar\beta^2 Q(E\cdot\omega)\int d^4p p_0 V''_{\beta p_0,\beta\mu}\delta(p^2-m^2),
 \end{aligned} \tag{A1}$$

where the second term in the first equality is vanishing. This can be seen by

$$\begin{aligned}
 &\tilde{\Omega}^{\mu\sigma}\partial_\mu(\beta u_\rho)\int d^4p p_\sigma p^\rho V''_{\beta p_0,\beta p_0}\delta(p^2-m^2) \\
 &= \tilde{\Omega}^{\mu\sigma}\partial_\mu(\beta u_\rho)u^\rho u_\sigma\int d^4p p_0^2 V''_{\beta p_0,\beta p_0}\delta(p^2-m^2) \\
 &\quad +\frac{1}{3}\tilde{\Omega}^{\mu\sigma}\partial_\mu(\beta u_\rho)\Delta_\sigma^\rho\int d^4p \bar{p}^2 V''_{\beta p_0,\beta p_0}\delta(p^2-m^2) \\
 &= 0,
 \end{aligned} \tag{A2}$$

where we have assumed that  $\beta$  is constant and used  $\Delta_\sigma^\rho = g_\sigma^\rho - u_\sigma u^\rho$ ,  $u^\rho\partial_\mu u_\rho = 0$  and

$$\tilde{\Omega}^{\mu\sigma}\partial_\mu(\beta u_\rho)\Delta_\sigma^\rho = \beta\tilde{\Omega}^{\mu\sigma}\partial_\mu u_\sigma = 2\beta\partial_\alpha\omega^\alpha = 0. \tag{A3}$$



Then we look at the second term in Eq. (15) which is related to electromagnetic field,

$$\begin{aligned}
I_F &= \hbar\beta Q \tilde{F}^{\mu\lambda} \int d^4p [-p_\lambda(\partial_\mu\mu) + (\partial_\mu u_\rho)p_\lambda p^\rho] V'_{\beta p_0} \delta'(p^2 - m^2) \\
&= \hbar\beta Q^2(E \cdot B) \int d^4p p_0 V'_{\beta p_0} \delta'(p^2 - m^2) \\
&\quad + \frac{1}{3} \hbar\beta Q \tilde{F}^{\mu\lambda} (\partial_\mu u_\lambda) \int d^4p \bar{p}^2 V'_{\beta p_0} \delta'(p^2 - m^2) \\
&= \hbar\beta Q^2(E \cdot B) \int d^4p p_0 V'_{\beta p_0} \delta'(p^2 - m^2) \\
&\quad + \frac{2}{3} \hbar\beta Q(E \cdot \omega) \int d^4p \bar{p}^2 V'_{\beta p_0} \delta'(p^2 - m^2) \\
&= \frac{1}{2} \hbar Q^2 \beta^2 (E \cdot B) \int d^4p V''_{\beta\mu, \beta p_0} \delta(p^2 - m^2) \\
&\quad + \hbar Q \beta (\omega \cdot E) \int d^4p \frac{1}{3p_0^2} \bar{p}^2 V'_{\beta p_0} \delta(p^2 - m^2) \\
&\quad - \hbar Q \beta^2 (\omega \cdot E) \int d^4p \frac{1}{3p_0} \bar{p}^2 V''_{\beta p_0, \beta p_0} \delta(p^2 - m^2), \tag{A4}
\end{aligned}$$

where we have used  $\delta'(x) = -\delta(x)/x$ ,  $\delta'(p^2 - m^2) \equiv d\delta(p^2 - m^2)/dp_0^2$ ,  $\partial_\mu\mu = -QE_\mu$ , and

$$\begin{aligned}
\partial_\mu V &= \beta [(\partial_\mu\mu)V'_{\beta\mu} + (\partial_\mu u_\rho)p^\rho V'_{\beta p_0}] \\
\frac{d}{dp_0}(\partial_\mu V) &= \frac{d}{dp_0} [\partial_\mu(\beta\mu)V'_{\beta\mu} + \partial_\mu(\beta u_\rho)p^\rho V'_{\beta p_0}] \\
&= \beta \partial_\mu(\beta\mu)V''_{\beta\mu, \beta p_0} + \partial_\mu(\beta u_\rho)u^\rho V'_{\beta p_0} + \beta \partial_\mu(\beta u_\rho)p^\rho V''_{\beta p_0, \beta p_0} \\
&= \beta^2 [(\partial_\mu\mu)V''_{\beta\mu, \beta p_0} + (\partial_\mu u_\rho)p^\rho V''_{\beta p_0, \beta p_0}] \\
(\partial_\mu u_\lambda)\tilde{F}^{\mu\lambda} &= \frac{1}{2} \epsilon^{\mu\lambda\alpha\beta} (\partial_\mu u_\lambda) F_{\alpha\beta} \\
&= \frac{1}{2} \epsilon^{\mu\lambda\alpha\beta} (\partial_\mu u_\lambda) (E_\alpha u_\beta - E_\beta u_\alpha + \epsilon_{\alpha\beta\rho\sigma} u^\rho B^\sigma) \\
&= 2\omega \cdot E + \frac{1}{2} \epsilon^{\mu\lambda\alpha\beta} \epsilon_{\alpha\beta\rho\sigma} (\partial_\mu u_\lambda) u^\rho B^\sigma \\
&= 2\omega \cdot E - (\partial_\rho u_\sigma - \partial_\sigma u_\rho) u^\rho B^\sigma \\
&= 2\omega \cdot E. \tag{A5}
\end{aligned}$$

We can add  $I_\omega$  and  $I_F$  from Eqs. (A1,A4) to obtain the right-hand side of Eq. (16), which we denote as  $I = I_\omega + I_F = I_{E \cdot B} + I_{E \cdot \omega}$  where  $I_{E \cdot B}$  and  $I_{E \cdot \omega}$  denote the  $E \cdot B$  and  $E \cdot \omega$  terms respectively. We now work on  $I_{E \cdot B}$  and  $I_{E \cdot \omega}$ . We now evaluate  $I_{E \cdot B}$  as,

$$\begin{aligned}
I_{E \cdot B} &= -\frac{1}{2} \hbar Q^2 \beta^2 (E \cdot B) \int d^4p V''_{\beta p_0, \beta p_0} \delta(p^2 - m^2) \\
&= \hbar Q^2 \beta^2 (E \cdot B) \int \frac{d^3p}{(2\pi)^3} \frac{1}{E_p} \\
&\quad \frac{d}{d(\beta E_p)} \{f_{\text{FD}}(E_p - \mu)[1 - f_{\text{FD}}(E_p - \mu)] + f_{\text{FD}}(E_p + \mu)[1 - f_{\text{FD}}(E_p + \mu)]\} \\
&= -\hbar Q^2 \beta (E \cdot B) \int \frac{d^3p}{(2\pi)^3} \frac{1}{E_p^2 - m^2} \\
&\quad \{f_{\text{FD}}(E_p - \mu)[1 - f_{\text{FD}}(E_p - \mu)] + f_{\text{FD}}(E_p + \mu)[1 - f_{\text{FD}}(E_p + \mu)]\} \\
&= -\frac{1}{2\pi^2} \hbar Q^2 (E \cdot B) C_1(\beta m, \beta\mu), \tag{A6}
\end{aligned}$$

where we have used  $d^3p = dE_p E_p \sqrt{E_p^2 - m^2}$  and  $C_1(\beta m, \beta \mu)$  is given in Eq. (18). Then the result of  $I_{E \cdot \omega}$  is

$$\begin{aligned}
I_{E \cdot \omega} &= \frac{1}{2} \hbar \beta Q(E \cdot \omega) \int d^4p \left[ \frac{2\bar{p}^2}{3p_0^2} V'_{\beta p_0} - \beta \left( p_0 + \frac{2\bar{p}^2}{3p_0} \right) V''_{\beta p_0, \beta p_0} \right] \delta(p^2 - m^2) \\
&= \frac{2}{3} \hbar \beta Q(E \cdot \omega) \int \frac{d^3p}{(2\pi)^3 E_p} \left( 1 - \frac{m^2}{E_p^2} \right) \\
&\quad \{ f_{\text{FD}}(E_p - \mu) [1 - f_{\text{FD}}(E_p - \mu)] - f_{\text{FD}}(E_p + \mu) [1 - f_{\text{FD}}(E_p + \mu)] \} \\
&\quad + \frac{1}{3} \hbar \beta^2 Q(E \cdot \omega) \int \frac{d^3p}{(2\pi)^3} \left( 1 + \frac{2m^2}{E_p^2} \right) \\
&\quad \{ [f_{\text{FD}}(E_p - \mu)] [1 - f_{\text{FD}}(E_p - \mu)] [2f_{\text{FD}}(E_p - \mu) - 1] \\
&\quad - [f_{\text{FD}}(E_p + \mu)] [1 - f_{\text{FD}}(E_p + \mu)] [2f_{\text{FD}}(E_p + \mu) - 1] \}. \tag{A7}
\end{aligned}$$

The last term in Eq. (A7) can be further simplified by using the formula and the integral by part for  $E_p$ ,

$$\begin{aligned}
\frac{d}{d(\beta x)} \{ f_{\text{FD}}(x) [1 - f_{\text{FD}}(x)] \} &= f_{\text{FD}}(x) [1 - f_{\text{FD}}(x)] [2f_{\text{FD}}(x) - 1], \\
\frac{d}{dE_p} \left[ E_p \sqrt{E_p^2 - m^2} \left( 1 + \frac{2m^2}{E_p^2} \right) \right] &= E_p \sqrt{E_p^2 - m^2} \left[ \left( \frac{1}{E_p} + \frac{E_p}{E_p^2 - m^2} \right) \left( 1 + \frac{2m^2}{E_p^2} \right) - 4 \frac{m^2}{E_p^3} \right]. \tag{A8}
\end{aligned}$$

Then Eq. (A7) can be further simplified as

$$\begin{aligned}
I_{E \cdot \omega} &= -\hbar \beta Q(E \cdot \omega) \int \frac{d^3p}{(2\pi)^3 E_p} \cdot \frac{m^2}{E_p^2 - m^2} \\
&\quad \{ f_{\text{FD}}(E_p - \mu) [1 - f_{\text{FD}}(E_p - \mu)] - f_{\text{FD}}(E_p + \mu) [1 - f_{\text{FD}}(E_p + \mu)] \} \\
&= -\frac{m^2}{2\pi^2} \hbar Q \beta (E \cdot \omega) C_2(\beta m, \beta \mu), \tag{A9}
\end{aligned}$$

where  $C_2(\beta m, \beta \mu)$  is given by Eq. (18). The right-hand side of Eq. (15) is just  $I_{E \cdot B} + I_{E \cdot \omega}$ , where  $I_{E \cdot B}$  and  $I_{E \cdot \omega}$  are in Eqs. (A6, A9) respectively, which finally gives Eq. (17).

## Appendix B: Small mass expansion of $C_1(\beta m, \beta \mu)$ and $C_2(\beta m, \beta \mu)$

In this appendix, we expand  $C_1(\beta m, \beta \mu)$  and  $C_2(\beta m, \beta \mu)$  in small  $\beta m$ . For simplicity of notations, we use new variables  $\bar{\mu} \equiv \beta \mu$ ,  $\bar{m} \equiv \beta m$  and define two dimensionless functions

$$\begin{aligned}
h(\bar{m}, \bar{\mu}) &= \int_0^\infty dx \frac{1}{e^{\sqrt{x^2 + \bar{m}^2} + \bar{\mu}} + 1}, \\
g(\bar{m}, \bar{\mu}) &= \int_0^\infty dx \frac{1}{\sqrt{x^2 + \bar{m}^2}} \frac{1}{e^{\sqrt{x^2 + \bar{m}^2} + \bar{\mu}} + 1}. \tag{B1}
\end{aligned}$$

We can express  $C_1(\bar{m}, \bar{\mu})$  and  $C_2(\bar{m}, \bar{\mu})$  in terms of  $h(\bar{m}, \bar{\mu})$  and  $g(\bar{m}, \bar{\mu})$ ,

$$\begin{aligned}
C_1(\bar{m}, \bar{\mu}) &= \frac{d}{d\bar{\mu}} [h(\bar{m}, -\bar{\mu}) - h(\bar{m}, \bar{\mu})], \\
C_2(\bar{m}, \bar{\mu}) &= \frac{d}{d\bar{\mu}} [g(\bar{m}, -\bar{\mu}) + g(\bar{m}, \bar{\mu})]. \tag{B2}
\end{aligned}$$

The Fermi-Dirac distribution can be expanded as

$$\frac{1}{e^{\sqrt{x^2 + \bar{m}^2} + \bar{\mu}} + 1} = \sum_{n=1}^{\infty} (-1)^{n+1} e^{-n(\sqrt{x^2 + \bar{m}^2} + \bar{\mu})}, \tag{B3}$$

and then we have

$$\begin{aligned}
h(\bar{m}, \bar{\mu}) &= \sum_{n=1}^{\infty} (-1)^{n+1} e^{-n\bar{\mu}} \int_0^{\infty} dx e^{-n\sqrt{x^2 + \bar{m}^2}} \\
&= \bar{m} \sum_{n=1}^{\infty} (-1)^{n+1} e^{-n\bar{\mu}} K_1(n\bar{m}), \\
g(\bar{m}, \bar{\mu}) &= \sum_{n=1}^{\infty} (-1)^{n+1} e^{-n\bar{\mu}} \int_0^{\infty} dx \frac{1}{\sqrt{x^2 + \bar{m}^2}} e^{-n\sqrt{x^2 + \bar{m}^2}} \\
&= \sum_{n=1}^{\infty} (-1)^{n+1} e^{-n\bar{\mu}} K_0(n\bar{m}),
\end{aligned} \tag{B4}$$

where  $K_0(x)$  and  $K_1(x)$  are modified Bessel functions of the second kind whose expansion form are

$$\begin{aligned}
K_0(x) &= -(\ln \frac{x}{2} + \gamma) \sum_{k=0}^{\infty} \frac{1}{(k!)^2} \left(\frac{x}{2}\right)^{2k} + \sum_{k=1}^{\infty} \frac{1}{(k!)^2} \left(1 + \frac{1}{2} + \dots + \frac{1}{k}\right) \left(\frac{x}{2}\right)^{2k}, \\
K_1(x) &= \frac{1}{x} + \left(\ln \frac{x}{2} + \gamma\right) \sum_{k=0}^{\infty} \frac{1}{k!(k+1)!} \left(\frac{x}{2}\right)^{2k+1} \\
&\quad - \frac{1}{2} \sum_{k=0}^{\infty} \frac{1}{k!(k+1)!} \left(1 + \frac{1}{2} + \dots + \frac{1}{k+1}\right) \left(\frac{x}{2}\right)^{2k+1} \\
&\quad - \frac{1}{2} \sum_{k=0}^{\infty} \frac{1}{(k+1)!(k+2)!} \left(1 + \frac{1}{2} + \dots + \frac{1}{k+1}\right) \left(\frac{x}{2}\right)^{2k+3},
\end{aligned} \tag{B5}$$

with the Euler constant  $\gamma = 0.577 \dots$ .

Then one can verify that  $h(\bar{m}, \bar{\mu})$  can be cast into the following form

$$\begin{aligned}
h(\bar{m}, \bar{\mu}) &= \phi(\bar{\mu}) - \bar{m} \sum_{k=0}^{\infty} \frac{1}{k!(k+1)!} \left(\frac{\bar{m}}{2} \frac{d}{d\bar{\mu}}\right)^{2k+1} \left[ \chi(\bar{\mu}) + \frac{1}{e^{\bar{\mu}} + 1} \ln \frac{\bar{m} e^{\gamma}}{2} \right] \\
&\quad + \frac{\bar{m}}{2} \sum_{k=0}^{\infty} \frac{1}{k!(k+1)!} \left(1 + \frac{1}{2} + \dots + \frac{1}{k+1}\right) \left(\frac{\bar{m}}{2} \frac{d}{d\bar{\mu}}\right)^{2k+1} \frac{1}{e^{\bar{\mu}} + 1} \\
&\quad + \frac{\bar{m}}{2} \sum_{k=0}^{\infty} \frac{1}{(k+1)!(k+2)!} \left(1 + \frac{1}{2} + \dots + \frac{1}{k+1}\right) \left(\frac{\bar{m}}{2} \frac{d}{d\bar{\mu}}\right)^{2k+3} \frac{1}{e^{\bar{\mu}} + 1},
\end{aligned} \tag{B6}$$

where we have defined two series

$$\begin{aligned}
\chi(\bar{\mu}) &= \sum_{n=1}^{\infty} (-1)^{n+1} e^{-n\bar{\mu}} \ln n, \\
\phi(\bar{\mu}) &= \sum_{n=1}^{\infty} (-1)^{n+1} e^{-n\bar{\mu}} \frac{1}{n}.
\end{aligned} \tag{B7}$$

Obviously we have following relation for  $\phi(\bar{\mu})$ ,

$$\begin{aligned}
\frac{d}{d\bar{\mu}} \phi(\bar{\mu}) &= - \sum_{n=1}^{\infty} (-1)^{n+1} e^{-n\bar{\mu}} = - \frac{1}{e^{\bar{\mu}} + 1}, \\
\frac{d}{d\bar{\mu}} \phi(-\bar{\mu}) &= - \frac{d}{d(-\bar{\mu})} \phi(-\bar{\mu}) = \frac{e^{\bar{\mu}}}{e^{\bar{\mu}} + 1}.
\end{aligned} \tag{B8}$$

Then we obtain

$$\begin{aligned}
h(\bar{m}, -\bar{\mu}) - h(\bar{m}, \bar{\mu}) &= \phi(-\bar{\mu}) - \phi(\bar{\mu}) + \bar{m} I_1 \left( \bar{m} \frac{d}{d\bar{\mu}} \right) [\chi(\bar{\mu}) + \chi(-\bar{\mu})], \\
\frac{d}{d\bar{\mu}} [h(\bar{m}, -\bar{\mu}) - h(\bar{m}, \bar{\mu})] &= 1 + \bar{m} I_1 \left( \bar{m} \frac{d}{d\bar{\mu}} \right) \frac{d}{d\bar{\mu}} [\chi(\bar{\mu}) + \chi(-\bar{\mu})],
\end{aligned} \tag{B9}$$

with the modified Bessel function  $I_1(x)$  of the first kind is given by

$$I_1(x) = \sum_{k=0}^{\infty} \frac{1}{k!(k+1)!} \left(\frac{x}{2}\right)^{2k+1}. \quad (\text{B10})$$

We can also treat  $g(\bar{m}, \bar{\mu})$  in the same way

$$\begin{aligned} g(\bar{m}, \bar{\mu}) &= - \sum_{k=0}^{\infty} \frac{1}{(k!)^2} \left(\frac{\bar{m}}{2} \frac{d}{d\bar{\mu}}\right)^{2k} \left[ \chi(\bar{\mu}) + \frac{1}{e^{\bar{\mu}} + 1} \ln \frac{\bar{m}e^{\gamma}}{2} \right] \\ &\quad + \sum_{k=1}^{\infty} \frac{1}{(k!)^2} \left(1 + \frac{1}{2} + \cdots + \frac{1}{k}\right) \left(\frac{\bar{m}}{2} \frac{d}{d\bar{\mu}}\right)^{2k} \frac{1}{e^{\bar{\mu}} + 1}, \end{aligned} \quad (\text{B11})$$

and we obtain

$$\begin{aligned} g(\bar{m}, -\bar{\mu}) + g(\bar{m}, \bar{\mu}) &= -\ln \frac{\bar{m}e^{\gamma}}{2} - I_0\left(\bar{m} \frac{d}{d\bar{\mu}}\right) [\chi(\bar{\mu}) + \chi(-\bar{\mu})], \\ \frac{d}{d\bar{\mu}} [g(\bar{m}, -\bar{\mu}) + g(\bar{m}, \bar{\mu})] &= -I_0\left(\bar{m} \frac{d}{d\bar{\mu}}\right) \frac{d}{d\bar{\mu}} [\chi(\bar{\mu}) + \chi(-\bar{\mu})], \end{aligned} \quad (\text{B12})$$

with the modified Bessel function  $I_0(x)$  of the first kind is given by

$$I_0(x) = \sum_{k=0}^{\infty} \frac{1}{(k!)^2} \left(\frac{x}{2}\right)^{2k}. \quad (\text{B13})$$

Finally we obtain the expressions for  $C_1(\bar{m}, \bar{\mu})$  and  $C_2(\bar{m}, \bar{\mu})$ ,

$$\begin{aligned} C_1(\bar{m}, \bar{\mu}) &= 1 + \bar{m} I_1\left(\bar{m} \frac{d}{d\bar{\mu}}\right) \frac{d}{d\bar{\mu}} [\chi(\bar{\mu}) + \chi(-\bar{\mu})], \\ C_2(\bar{m}, \bar{\mu}) &= -I_0\left(\bar{m} \frac{d}{d\bar{\mu}}\right) \frac{d}{d\bar{\mu}} [\chi(\bar{\mu}) + \chi(-\bar{\mu})]. \end{aligned} \quad (\text{B14})$$

The series  $\chi(\bar{\mu})$  can also be written in the form,

$$\begin{aligned} \chi(\bar{\mu}) &= \sum_{n=1}^{\infty} (-1)^{n+1} e^{-n\bar{\mu}} \ln n = \left[ \frac{\partial}{\partial y} \sum_{n=1}^{\infty} \frac{(-e^{-\bar{\mu}})^n}{n^y} \right]_{y=0} \\ &= \left[ \frac{\partial}{\partial y} \text{Li}(y, -e^{-\bar{\mu}}) \right]_{y=0}, \end{aligned} \quad (\text{B15})$$

where  $\text{Li}(y, z)$  is the polylogarithm function,

$$\text{Li}(y, z) = \sum_{n=1}^{\infty} \frac{z^n}{n^y}, \quad |z| < 1. \quad (\text{B16})$$

The analytic continuation of the polylogarithm beyond the circle of convergence  $|z| < 1$  can be furnished by following integral representation,

$$\text{Li}(y, z) = \frac{1}{\Gamma(y)} \int_0^{\infty} \frac{t^{y-1}}{z^{-1}e^t - 1} dt, \quad \text{Re}(y) > 0, \quad z \in \mathbb{C}/[1, \infty), \quad (\text{B17})$$

with the Gamma function  $\Gamma(y)$  defined by

$$\Gamma(y) = \int_0^{\infty} e^{-t} t^{y-1} dt, \quad \text{Re}(y) > 0. \quad (\text{B18})$$

The asymptotic behaviors of  $I_0(x)$  and  $I_1(x)$  at  $x \rightarrow 0$  are  $I_0(x) \approx 1$  and  $I_1(x) \approx x/2$ . So at  $\bar{m} \rightarrow 0$ , we obtain the asymptotic values of  $C_1(\bar{m}, \bar{\mu})$  and  $C_2(\bar{m}, \bar{\mu})$ ,

$$\begin{aligned} C_1(\bar{m}, \bar{\mu}) &\approx 1 + \frac{\bar{m}^2}{2} \frac{d^2}{d\bar{\mu}^2} \frac{d}{dy} [\text{Li}(y, -e^{-\bar{\mu}}) + \text{Li}(y, -e^{\bar{\mu}})]_{y=0}, \\ C_2(\bar{m}, \bar{\mu}) &\approx -\frac{d}{d\bar{\mu}} \frac{d}{dy} [\text{Li}(y, -e^{-\bar{\mu}}) + \text{Li}(y, -e^{\bar{\mu}})]_{y=0}. \end{aligned} \quad (\text{B19})$$

### Appendix C: Derivation of Eq. (20)

In this appendix, we give a detailed derivation of Eq. (20). From the definition of the Wigner function (3) and that of the axial vector component, we obtain

$$\begin{aligned}
QF_{\mu\nu} \int d^4p \partial_p^\nu \mathcal{A}^\mu &= QF_{\mu\nu} \int d^4p \partial_p^\nu \int \frac{d^4y}{(2\pi)^4} e^{-ip \cdot y} \\
&\quad \times \left\langle \bar{\psi}(x + \frac{1}{2}y) \gamma^\mu \gamma^5 PU(G, x + \frac{1}{2}y, x - \frac{1}{2}y) \psi(x - \frac{1}{2}y) \right\rangle \\
&= QF_{\mu\nu} \int d^4y (-iy^\nu) \delta^{(4)}(y) \\
&\quad \times \left\langle \bar{\psi}(x + \frac{1}{2}y) \gamma^\mu \gamma^5 PU(G, x + \frac{1}{2}y, x - \frac{1}{2}y) \psi(x - \frac{1}{2}y) \right\rangle \\
&= -iQF_{\mu\nu} \left( \lim_{y \rightarrow 0} y^\nu \left\langle \bar{\psi}(x + \frac{y}{2}) \gamma^\mu \gamma^5 \psi(x - \frac{y}{2}) \right\rangle \right) \\
&= -\frac{Q^2}{8\pi^2} F_{\mu\nu} \tilde{F}^{\mu\nu} \\
&= -\frac{Q^2}{2\pi^2} B \cdot E
\end{aligned} \tag{C1}$$

where we have used  $4B \cdot E = F_{\mu\nu} \tilde{F}^{\mu\nu}$ ,  $\lim_{y \rightarrow 0} PU(G, x + \frac{1}{2}y, x - \frac{1}{2}y) = 1$  and [58]

$$\begin{aligned}
\lim_{y \rightarrow 0} \bar{\psi}(x + \frac{y}{2}) \gamma^\mu \gamma^5 \psi(x - \frac{y}{2}) &= -\frac{i}{4\pi^2} Q \epsilon^{\alpha\beta\mu\rho} F_{\alpha\beta} \lim_{y \rightarrow 0} \frac{y_\rho}{y^2} \\
\lim_{y \rightarrow 0} \frac{y^\nu y^\rho}{y^2} &= \frac{1}{4} g^{\nu\rho}
\end{aligned} \tag{C2}$$

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